Correlative signatures of heavy Majorana neutrinos and doubly-charged Higgs bosons at the Large Hadron Collider

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Abstract

We explore an intriguing possibility to test the type-II seesaw mechanism of neutrino mass generation at the Large Hadron Collider (LHC). We show that the lepton-number-violating signatures of heavy Majorana neutrinos N_i (for $i = 1, 2, 3$) and doubly-charged Higgs bosons $H^{\pm\pm}$ can be closely correlated with each other in a class of TeV-scale type-II seesaw models. Taking the minimal version of such models for example, we calculate the cross sections of $pp \to l^\pm_\alpha l^\pm_\beta X$ processes mediated separately by N_1 and $H^{\pm\pm}$, and illustrate their nontrivial correlation at the LHC.

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I. INTRODUCTION

The running of the Large Hadron Collider (LHC) in the coming years will shed light on several fundamental problems in the Standard Model (SM) and explore possible new physics beyond the SM [1]. If new physics exists at the TeV scale and is responsible for the electroweak symmetry breaking, it may also be responsible for the origin of neutrino masses. The latter is a kind of new physics which has been firmly established by a number of neutrino oscillation experiments in the past decade [2]. Therefore, it is extremely interesting to see whether some deep understanding of the neutrino mass generation and lepton number (flavor) violation can be achieved at the energy frontier set by the LHC.

A natural and attractive possibility of generating tiny neutrino masses is to extend the SM by introducing three heavy right-handed Majorana neutrinos [3] and (or) one Higgs triplet [4]. The gauge-invariant Lagrangian relevant to lepton masses can then be written as

$$-\mathcal{L}_{\rm lepton} = \overline{l_{\rm L}} Y_l H E_{\rm R} + \overline{l_{\rm L}} Y_\nu \tilde{H} N_{\rm R} + \frac{1}{2} \overline{N^c_{\rm R}} M_{\rm R} N_{\rm R} + \frac{1}{2} \overline{l_{\rm L}} Y_\Delta \Delta i\sigma_2 l^c_{\rm L} + {\rm h.c.} \; , \tag{1}$$

where Δ is the Higgs triplet, and $M_{\rm R}$ is the mass matrix of right-handed Majorana neutrinos. After the spontaneous gauge symmetry breaking, one obtains the mass matrices $M_l = Y_l v/\sqrt{2}$, $M_{\rm D} = Y_\nu v/\sqrt{2}$ and $M_{\rm L} = Y_\Delta v_\Delta$, where $\langle H \rangle \equiv v/\sqrt{2}$ and $\langle \Delta \rangle \equiv v_\Delta$ are the vacuum expectation values (vev's) of the neutral components of H and Δ, respectively. In the leading-order approximation, the effective mass matrix for three light neutrinos is given by $M_\nu \approx M_{\rm L} - M_{\rm D} M^{-1}_{\rm R} M^T_{\rm D}$. This is the so-called type-II seesaw mechanism. If the Higgs triplet Δ is absent, the small mass scale of M_ν can be just attributed to the large mass scale of $M_{\rm R}$ (type-I seesaw [3]). In the absence of heavy right-handed Majorana neutrinos, the small mass scale of M_ν implies that the mass scale of $M_{\rm L}$ must be equally small (triplet seesaw [4]). Another interesting case is that both terms of M_ν are important and their significant cancellation gives rise to small neutrino masses [5]. Direct tests of such neutrino mass models can be done at the LHC, provided they work at the TeV scale and predict appreciable collider signatures.

The search for the Higgs triplet and heavy Majorana neutrinos at the Tevatron and LHC has recently attracted a lot of attention. For example, a model-independent analysis has been done in Ref. [6] to probe the same-sign dilepton events induced by heavy Majorana neutrinos via the most promising channel $q\bar{q}' \to \mu^\pm N_i \to \mu^\pm \mu^\pm W^\mp$. These events signify the lepton number violation and serve for a clean collider signature of new physics beyond

the SM [7]. In *realistic* type-I seesaw models, however, the observability of N_i requires $\mathcal{O}(M_{\rm R}) \lesssim 1$ TeV and $\mathcal{O}(M_{\rm D}/M_{\rm R}) \gtrsim 10^{-3}$ together with an unnatural cancellation condition $M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T \approx \mathbf{0}$ [8, 9]. In a triplet seesaw model, the production of the doubly-charged Higgs bosons $H^{\pm\pm}$ only depends on their masses and has nothing to do with the Yukawa coupling Y_Δ. It is therefore possible to discover $H^{\pm\pm}$ of $\mathcal{O}(\lesssim 1)$ TeV at the LHC by detecting $H^{\pm\pm} \to l^\pm_\alpha l^\pm_\beta$ decays [10, 11, 12, 13].

Different from those previous works, this paper aims to investigate possible correlation between the collider signatures of heavy Majorana neutrinos N_i and doubly-charged Higgs bosons $H^{\pm\pm}$ in a class of realistic type-II seesaw models, in which the smallness of M_ν is ascribed to a significant but incomplete cancellation between large $M_{\rm L}$ and $M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ terms. We shall see that the phenomenology of such a type-II seesaw model is much richer than that of a type-I seesaw model or a triplet seesaw model at the TeV scale. In particular, the non-unitarity of the light neutrino mixing matrix can be closely related to the lepton-number-violating signals of both N_i and $H^{\pm\pm}$ at the LHC. Taking the minimal version of the type-II seesaw models (with only one heavy Majorana neutrino N_1 in addition to the Higgs triplet Δ) for example, we calculate the cross sections of $pp \to l^\pm_\alpha l^\pm_\beta X$ mediated separately by N_1 and $H^{\pm\pm}$, and illustrate their nontrivial correlation at the LHC. We arrive at some interesting and encouraging conclusions.

II. THE MODEL

If an $SU(2)_{\rm L}$ Higgs triplet is introduced into the SM [4], the gauge-invariant potential can be written as $V(H,\Delta) = V_{\rm SM}(H) + \delta V$, where $V_{\rm SM}(H) = -\mu^2 H^\dagger H + \lambda \left(H^\dagger H\right)^2$ with $H \equiv (\varphi^+, \varphi^0)^T$ being the SM Higgs doublet, and

$$\delta V = \frac{1}{2} M_\Delta^2 {\rm Tr}\left(\Delta^\dagger \Delta\right) - \left[\lambda_\Delta M_\Delta H^T i\sigma_2 \Delta H + {\rm h.c.}\right] \qquad (2)$$

with Δ being defined as

$$\Delta \equiv \left(\begin{matrix} \xi^- & -\sqrt{2}\,\xi^0 \\ \sqrt{2}\,\xi^{--} & -\xi^- \end{matrix} \right) . \qquad (3)$$

When the neutral components of H and Δ acquire their vev's v and v_Δ, respectively, the electroweak gauge symmetry is spontaneously broken. The minimum of $V(H,\Delta)$ can be achieved at $v = \mu/(\lambda - 2\lambda_\Delta^2)^{1/2}$ and $v_\Delta = \lambda_\Delta v^2/M_\Delta$, where the dimensionless parameter λ_Δ has been assumed to be real. Note that v_Δ may affect the masses of $W^\pm$ and Z^0 in

such a way that $\rho \equiv M_W^2/(M_Z^2 \cos^2\theta_{\rm W}) = (v^2 + 2v_\Delta^2)/(v^2 + 4v_\Delta^2)$ holds. By using current data on the ρ-parameter [2], we get $\kappa \equiv \sqrt{2}\ v_\Delta/v < 0.01$ and $v_\Delta < 2.5$ GeV. There are totally seven physical Higgs bosons in this model: doubly-charged H^{++} and H^{--}, singly-charged H^+ and H^-, neutral A^0 (CP-odd), and neutral h^0 and H^0 (CP-even), where h^0 is the SM-like Higgs boson. Except for $M_{h^0}^2 \approx 2\mu^2$, we get a quasi-degenerate mass spectrum for other scalars: $M_{H^{\pm\pm}}^2 = M_\Delta^2 \approx M_{H^0}^2$, $M_{H^\pm}^2 = M_\Delta^2(1+\kappa^2)$, and $M_{A^0}^2 = M_\Delta^2(1+2\kappa^2)$. As a consequence, the decay channels $H^{\pm\pm} \to W^\pm H^\pm$ and $H^{\pm\pm} \to H^\pm H^\pm$ are kinematically forbidden. The production of $H^{\pm\pm}$ at the LHC is mainly through $q\bar{q} \to \gamma^*, Z^* \to H^{++}H^{--}$ and $q\bar{q}' \to W^* \to H^{\pm\pm}H^\mp$ processes.

On the neutrino side, we are left with $M_{\rm L}$, $M_{\rm D}$ and $M_{\rm R}$ from Eq. (1) after the electroweak symmetry breaking. They form a symmetric 6×6 matrix $\mathcal{M}$, which can be diagonalized by a unitary transformation $\mathcal{U}^\dagger \mathcal{M} \mathcal{U}^* = \widehat{\mathcal{M}}$. Explicitly,

$$\begin{pmatrix} V & R \\ S & U \end{pmatrix}^\dagger \begin{pmatrix} M_{\rm L} & M_{\rm D} \\ M_{\rm D}^T & M_{\rm R} \end{pmatrix} \begin{pmatrix} V & R \\ S & U \end{pmatrix}^* = \begin{pmatrix} \widehat{m} & \mathbf{0} \\ \mathbf{0} & \widehat{M} \end{pmatrix}, \tag{4}$$

$\widehat{m} = {\rm Diag}\{m_1, m_2, m_3\}$ and $\widehat{M} = {\rm Diag}\{M_1, M_2, M_3\}$ with m_i and M_i (for $i = 1,2,3$) being the light and heavy Majorana neutrino masses, respectively. In the mass basis, the leptonic charged-current interactions can be written as

$$-\mathcal{L}_{\rm cc} = \frac{g}{\sqrt{2}} \left[\overline{e_{\rm L}} V \gamma^\mu \nu_{\rm L} W_\mu^- + \overline{e_{\rm L}} R \gamma^\mu N_{\rm L} W_\mu^-\right] + {\rm h.c.}\ . \tag{5}$$

Note that $VV^\dagger + RR^\dagger = \mathbf{1}$ holds due to the unitarity of $\mathcal{U}$, and thus the neutrino mixing matrix V itself must be non-unitary [14]. The unitarity violation of V is characterized by R, which is responsible for the production and decays of heavy Majorana neutrinos N_i. In the leading-order approximation, the effective mass matrix of three light neutrinos is given by $M_\nu \approx M_{\rm L} - M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$. Either $M_{\rm L}$ or $M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ may dominate M_ν, but here we focus on the third possibility: the smallness of M_ν arises from a significant cancellation between $M_{\rm L}$ and $M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ in the case of $\mathcal{O}(M_\nu) \ll \mathcal{O}(M_{\rm L}) \sim \mathcal{O}(M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T)$. We admit that a substantial fine-tuning at the level of $\mathcal{O}(10^{-10})$, which is comparable with the degree of fine-tuning for the *structural cancellation* of $M_{\rm D}$ and $M_{\rm R}$ in the TeV-scale type-I seesaw models [8, 9], has to be required in order to realize this kind of *global cancellation* [1]. Let us reiterate the key points of our model in the following:

[1] In our model, $\mathcal{O}(M_{\rm L}) \sim \mathcal{O}(M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T) \sim 1$ GeV and $\mathcal{O}(M_\nu) \sim 0.1$ eV hold and thus the cancellation can be achieved with the precision of $\mathcal{O}(10^{-10})$. It is possible to realize such a fine cancellation in a more

- We assume that both M_i and M_Δ are of $\mathcal{O}(1)$ TeV. Then the production of $H^{\pm\pm}$ at the LHC is guaranteed, and their lepton-number-violating signatures will probe the Higgs triplet sector of the type-II seesaw mechanism. On the other hand, $\mathcal{O}(M_{\rm D}/M_{\rm R}) \lesssim 10\%$ is possible as a result of $\mathcal{O}(M_{\rm R}) \sim 1$ TeV and $\mathcal{O}(M_{\rm D}) \lesssim \mathcal{O}(v)$, such that appreciable signatures of N_i can be achieved at the LHC.

- The small mass scale of M_ν implies that the relation $\mathcal{O}(M_{\rm L}) \sim \mathcal{O}(M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T)$ must hold. In other words, it is the significant but incomplete cancellation between $M_{\rm L}$ and $M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ terms that results in the non-vanishing but tiny masses for three light neutrinos [2].

In this spirit, $M_{\rm L}$ can be reconstructed via the excellent approximation $M_{\rm L} = V\widehat{m}V^T + R\widehat{M}R^T \approx R\widehat{M}R^T$. The elements of the Yukawa coupling matrix Y_Δ are then given by

$$(Y_\Delta)_{\alpha\beta} = \frac{(M_{\rm L})_{\alpha\beta}}{v_\Delta} \approx \sum_{i=1}^{3} \frac{R_{\alpha i} R_{\beta i} M_i}{v_\Delta} \; , \tag{6}$$

where the subscripts α and β run over e, μ and τ. This result implies that the leptonic decays of $H^{\pm\pm}$ depend on both R and M_i, which actually determine the production and decays of N_i. Thus we have established an interesting correlation between the doubly-charged Higgs bosons and the heavy Majorana neutrinos. To observe the correlative signatures of $H^{\pm\pm}$ and N_i at the LHC will serve for a direct test of our type-II seesaw model.

We shall subsequently consider the minimal version of the type-II seesaw models [15], in which there is only one heavy Majorana neutrino N_1 together with the Higgs triplet Δ, to illustrate how the collider signatures of N_1 and $H^{\pm\pm}$ are correlated with each other. In this simple but instructive case, the decay rates of $H^{\pm\pm}$ are given by

$$\Gamma(H^{\pm\pm} \to l^\pm_\alpha l^\pm_\beta) \; = \; \frac{M_1^2 M_\Delta}{8\pi(1+\delta_{\alpha\beta}) v_\Delta^2} |R_{\alpha 1}|^2 |R_{\beta 1}|^2 \; , \tag{7}$$

or less natural way by imposing a certain flavor symmetry on $M_{\rm L}$, $M_{\rm D}$ and $M_{\rm R}$ and introducing slight perturbations to them [5]. In doing so, however, one has to be very cautious and should take into account the effects of radiative corrections to two mass terms [8, 9].

[2] Note that radiative corrections to M_ν can be expressed as $\delta M_\nu \approx -(M_{\rm L}\delta_{\rm L} - M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T \delta_{\rm R})$, where $|\delta_{\rm L}| \ll |\delta_{\rm R}| \lesssim 10^{-3}$ is numerically expected in our model. Hence Eq. (6) is a good approximation irrelevant to small radiative corrections. Although the magnitude of δM_ν is likely to be much larger than that of m_i (e.g., of $\mathcal{O}(1)$ MeV), it can in principle be suppressed via a kind of more subtle cancellation which may be accomplished by a slight modification of the relevant Yukawa couplings. A detailed analysis of radiative corrections to M_ν in the TeV-scale type-II seesaw model will be presented elsewhere.

which depend on M_1 and R in a very obvious way. Since R determines both the strength of non-unitarity of V and that of the charged-current interactions of N_1, it bridges a gap between neutrino physics and collider physics.

III. LHC SIGNATURES

We first look at the production of N_1 from proton-proton collisions at the LHC. The dominant channel is $pp \to l_\alpha^{+(-)} N_1 X \to l_\alpha^{+(-)} l_\beta^{+(-)} W^{-(+)} X$, in which N_1 is on-shell produced. Given the charged-current interactions in Eq. (5), the total cross section of this process reads

$$\sigma(pp \to l_\alpha^+ l_\beta^+ W^- X) \;\approx\; \sigma_N \cdot \frac{|R_{\alpha 1}|^2 |R_{\beta 1}|^2}{4 \sum\limits_\gamma |R_{\gamma 1}|^2} \;, \tag{8}$$

where $\sigma_N \equiv \sigma(pp \to l_\alpha^+ N_1 X)/|R_{\alpha 1}|^2$ and the narrow-width approximation has been used. Note that the reduced cross section σ_N is independent of any elements of R, and thus it can be computed by taking account of the parton distribution functions and the mass of N_1. If N_1 is much heavier than the gauge bosons and the SM-like Higgs boson, we have $\mathrm{Br}(N_1 \to l^+ W^-) \simeq \mathrm{Br}(N_1 \to \nu Z^0) \simeq \mathrm{Br}(N_1 \to \nu h^0) \approx 25\%$ to a very good degree of accuracy [8]. That is why there appears a factor $1/4$ on the right-hand side of Eq. (8).

We proceed to consider the production of $H^{\pm\pm}$ and their same-sign dilepton events at the LHC. In the narrow-width approximation, three relevant cross sections can be factorized as

$$\begin{aligned} \sigma(pp \to l_\alpha^+ l_\beta^+ H^- X) \;&=\; \sigma_H \cdot \mathrm{Br}(H^{++} \to l_\alpha^+ l_\beta^+) \;, \\ \sigma(pp \to l_\alpha^+ l_\beta^+ W^- X) \;&=\; \sigma_W \cdot \mathrm{Br}(H^{++} \to l_\alpha^+ l_\beta^+) \;, \\ \sigma(pp \to l_\alpha^+ l_\beta^+ H^{--} X) \;&=\; \sigma_{\rm pair} \cdot \mathrm{Br}(H^{++} \to l_\alpha^+ l_\beta^+) \;, \end{aligned} \tag{9}$$

where $\sigma_H \equiv \sigma(pp \to H^{++} H^- X)$, $\sigma_W \equiv \sigma(pp \to H^{++} W^- X)$, and $\sigma_{\rm pair} \equiv \sigma(pp \to H^{++} H^{--} X)$. The reduced cross sections σ_H, σ_W and $\sigma_{\rm pair}$ only depend on M_Δ and v_Δ, and they can be calculated in a way similar to the calculation of σ_N. To detect the lepton-number-violating signals of $H^{\pm\pm}$, we need to take account of their decay modes. Because of $M_{H^{\pm\pm}} \approx M_{H^\pm}$, only two decay modes are kinematically open: $H^{\pm\pm} \to l_\alpha^\pm l_\beta^\pm$ and $H^{\pm\pm} \to W^\pm W^\pm$ [11, 12]. The leptonic channel is expected to be dominant in our type-II seesaw model, and a detailed analysis of the branching ratios of $H^{\pm\pm}$ decays can be found in Ref. [16].

To specify the correlation between the signatures of N_1 and $H^{\pm\pm}$ at the LHC, let us parametrize the 3×1 complex matrix R in terms of three rotation angles and three phase angles [14]: $R = (\hat{s}^*_{14}, c_{14}\hat{s}^*_{24}, c_{14}c_{24}\hat{s}^*_{34})^T$, where $c_{ij} \equiv \cos\theta_{ij}$ and $\hat{s}_{ij} \equiv e^{i\delta_{ij}} s_{ij}$ with $s_{ij} \equiv \sin\theta_{ij}$ (for $ij = 14, 24, 34$). A global analysis of current neutrino oscillation data and precision electroweak data yields very stringent constraints on the non-unitarity of the neutrino mixing matrix V, which are equivalent to the constraints on R [17]: $s_{14}s_{24} \leq 7.0 \times 10^{-5}$ in addition to $s^2_{14}, s^2_{24}, s^2_{34} \leq 1.0 \times 10^{-2}$. On the other hand, the experimental upper bound on the neutrinoless double-beta decay requires $s^2_{14}(1 + M^2_1/M^2_\Delta)/M_1 < 5 \times 10^{-8}$ GeV^{-1}, where both the contributions of N_1 and $H^{\pm\pm}$ have been taken into account. Combining all these constraints, we may choose a typical and self-consistent parameter space of three mixing angles: $s_{14} \approx 0$, $s_{24} \in [0.01, 0.1]$ and $s_{34} \in [0.01, 0.1]$ [16]. The decay modes $H^{\pm\pm} \to e^\pm e^\pm$, $e^\pm\mu^\pm$ and $e^\pm\tau^\pm$ are therefore forbidden, while

$$\begin{aligned} \mathrm{Br}(H^{\pm\pm} \to \mu^\pm\mu^\pm) &\approx \frac{s^4_{24}}{\left(s^2_{24} + s^2_{34}\right)^2} \ , \\ \mathrm{Br}(H^{\pm\pm} \to \mu^\pm\tau^\pm) &\approx \frac{2s^2_{24}s^2_{34}}{\left(s^2_{24} + s^2_{34}\right)^2} \ , \end{aligned} \tag{10}$$

and $\mathrm{Br}(H^{\pm\pm} \to \tau^\pm\tau^\pm) \approx 1 - \mathrm{Br}(H^{\pm\pm} \to \mu^\pm\tau^\pm) - \mathrm{Br}(H^{\pm\pm} \to \mu^\pm\mu^\pm)$.

For each lepton-number-violating process $pp \to l^\pm_\alpha l^\pm_\beta X$ discussed above, its cross section is actually calculated in the following way:

$$\sigma = \sum_{a,b} \int \mathrm{d}x_1 \mathrm{d}x_2 F_{a/p}(x_1, Q^2) \cdot F_{b/p}(x_2, Q^2) \cdot \ \hat{\sigma}(ab \to l^\pm_\alpha l^\pm_\beta X) \ , \tag{11}$$

where F denotes the parton distribution function, $x_{1,2}$ is the energy fraction of the partons, Q is the factorization scale, and $\hat{\sigma}$ is the partonic cross section. For illustration, we fix $s_{14} = 0$, $s_{24} = s_{34} = 0.1$, $Q^2 = x_1x_2S$ with $\sqrt{S} = 14$ TeV, $v_\Delta = 1$ GeV and $v = 246$ GeV in our subsequent numerical calculations. We plot σ_N, σ_H, σ_W and $\sigma_{\rm pair}$ in FIG. 1 (a) by allowing M_1 and M_Δ to vary from 200 GeV to 2 TeV. Note again that these reduced cross sections are independent of s_{ij}. In FIG. 1 (a), we also show the required cross sections $\bar{\sigma}_{N(H)}$ for one $pp \to \mu^+\mu^+X$ event produced at the LHC with the integrated luminosity of 300 fb^{-1} in both the case of heavy Majorana neutrino N_1 (dashed line) and that of doubly-charged Higgs boson H^{++} (solid line). The cross section $\bar{\sigma}_{N(H)}$ is defined as follows

$$\bar{\sigma}_{N(H)} = \frac{1 \text{ event}}{300 \text{ fb}^{-1}} \, \frac{1}{B_{N(H)}}, \tag{12}$$

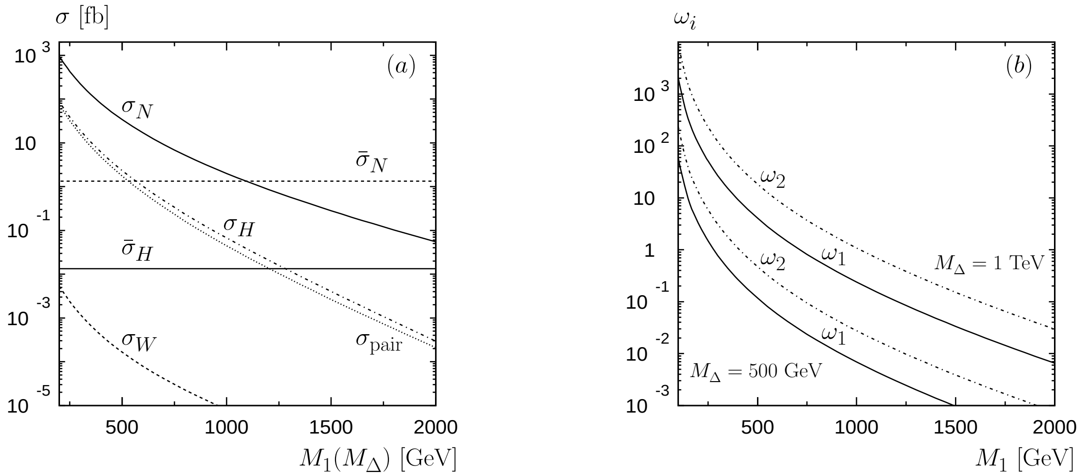

FIG. 1: (a) The reduced cross sections σ_N, σ_H, σ_W and $\sigma_{\rm pair}$ at the LHC, where M_1 and M_Δ are the masses of N_1 and $H^{\pm\pm}$, respectively. The horizontal dashed (solid) line corresponds to the cross section $\bar{\sigma}_{N(H)}$ for one event induced by heavy Majorana neutrinos (doubly-charged Higgs bosons) at the LHC with the integrated luminosity of 300 fb^{-1}. (b) The correlation between the lepton-number-violating signatures of N_1 and $H^{\pm\pm}$ at the LHC, where $s_{14} = 0$, $s_{24} = s_{34} = 0.1$ and $v_\Delta = 1$ GeV have typically been input.

where

$$B_N = \frac{|R_{\alpha 1}|^2 |R_{\beta 1}|^2}{4 \sum_\gamma |R_{\gamma 1}|^2}, \qquad B_H = {\rm Br}(H^{++} \to \mu^+ \mu^+). \tag{13}$$

It is obvious that the production rate of $pp \to H^{++} W^- X$ is remarkably smaller than those of other processes and can be neglected. One can find from our results that at the LHC, $H^{\pm\pm}$ (heavy Majorana neutrinos) should be observable with the integrated luminosity of 300 fb^{-1} in the $l^\pm l^\pm$ channel up to the mass of 1.2 TeV (1.1 TeV). Detailed analysis is not the aim of this paper and can be found in our following work. Furthermore, σ_H is larger than $\sigma_{\rm pair}$, and the ratio of σ_H to $\sigma_{\rm pair}$ lies in the range $1.1 \cdots 1.6$ for $M_\Delta \in [200~{\rm GeV}, 2~{\rm TeV}]$, which is consistent with the previous results [11, 13]. The correlation between the LHC signatures of N_1 and $H^{\pm\pm}$ becomes more transparent in

$$\begin{aligned} \omega_1 &\equiv \frac{\sigma(pp \to \mu^+ \mu^+ W^- X)|_{N_1}}{\sigma(pp \to \mu^+ \mu^+ H^- X)|_{H^{++}}} , \\ \omega_2 &\equiv \frac{\sigma(pp \to \mu^+ \mu^+ W^- X)|_{N_1}}{\sigma(pp \to \mu^+ \mu^+ H^{--} X)|_{H^{++}}} , \end{aligned} \tag{14}$$

which can approximate to $\omega_1 \approx \sigma_N (s_{24}^2 + s_{34}^2)/(4\sigma_H)$ and $\omega_2 \approx \sigma_N (s_{24}^2 + s_{34}^2)/(4\sigma_{\rm pair})$, respectively. Comparing between Eq. (8) and Eq. (10), we find that $\omega_{1,2}$ is universal for $\mu\mu$,

$\mu\tau$ and $\tau\tau$ modes. The changes of ω_1 and ω_2 with M_1 are illustrated in FIG. 1 (b), where $M_\Delta = 500$ GeV and 1 TeV have typically been input.

Finally, let us make some brief comments on the type-II seesaw models with two or three heavy Majorana neutrinos. If the masses of N_i are nearly degenerate, their production cross sections will be modified due to either constructive or destructive interference between different contributions. In particular, the CP-violating phases of R will enter the expressions of $\text{Br}(H^{\pm\pm} \rightarrow l^\pm_\alpha l^\pm_\beta)$ and influence the lepton-number-violating signatures of the doubly-charged Higgs bosons at the LHC. Because a generic type-II seesaw model contains many more free parameters even in the case of $M_\text{L} \approx R\widehat{M}R^T$, the study of its collider signatures will involve much more uncertainties. Such an analysis, which is important but beyond the scope of this paper, will be done elsewhere [16].

IV. SUMMARY

Motivated by the conjecture that new physics at the TeV scale might not only solve the naturalness problem of electroweak symmetry breaking but also be responsible for the lepton number (flavor) violation, we have explored an intriguing possibility to test the type-II seesaw mechanism of neutrino mass generation at the LHC. Our key point is that the mass scales of the $SU(2)_\text{L}$ Higgs triplet and heavy Majorana neutrinos are both of $\mathcal{O}(\lesssim 1)$ TeV, and the smallness of $M_\nu \approx M_\text{L} - M_\text{D} M^{-1}_\text{R} M^T_\text{D}$ is attributed to a significant cancellation between its M_L and $M_\text{D} M^{-1}_\text{R} M^T_\text{D}$ terms. This observation allows us to establish an interesting correlation between the lepton-number-violating signatures of heavy Majorana neutrinos N_i and doubly-charged Higgs bosons $H^{\pm\pm}$, but it has nothing to do with the mass spectrum of three light neutrinos (either a normal hierarchy or an inverted hierarchy). Taking the minimal version of the type-II seesaw models for example, we have calculated the cross sections of $pp \rightarrow l^\pm_\alpha l^\pm_\beta X$ mediated separately by N_1 and $H^{\pm\pm}$, and illustrated their nontrivial correlation at the LHC. Our results are quite encouraging, and our analysis can easily be extended to the more general cases of this class of TeV-scale type-II seesaw models.

We stress that it is extremely important to search for the correlative collider signatures of N_i and $H^{\pm\pm}$, so as to convincingly and unambiguously verify the type-II seesaw mechanism. In contrast, individual signatures of N_i or $H^{\pm\pm}$ are only possible to demonstrate the type-I or triplet seesaw mechanism. We also stress that the non-unitarity of the 3×3 neutrino mixing

matrix is intimately correlated with the LHC signatures of N_i. This kind of correlation, which bridges a gap between neutrino physics at low energies and collider physics at the TeV scale, deserves a lot of further investigation.

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